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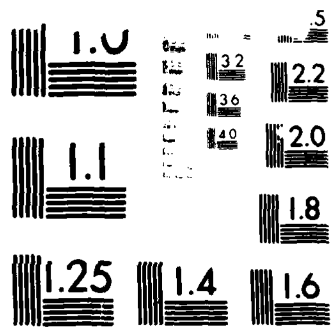
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Non-Maxwellian Free-Energy Generation in the Magnetotail Due to Chaotic Particle Motion

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It has previously been shown that the charged-particle motion in a magnetotail-like magnetic field is generally non-integrable and that the particle orbits can be classified into three distinct types: the bounded integrable orbits, unbounded stochastic orbits and unbounded transient orbits. The three distinct types of orbits occupy disjoint regions of the phase space. As a result, the different regions of the phase space retain the memory of the existing distribution function for different lengths of time. The concept of "differential memory" in single-particle distributions is discussed. Physical implications for the macroscopical properties of the magnetotail plasmas, and the possible generation of non-Maxwellian features of the distribution function are discussed.

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CONTENTS

NON-MAXWELLIAN FREE-ENERGY GENERATION IN THE MAGNETOTAIL DUE TO CHAOTIC PARTICLE MOTION

I. Introduction

The motion of charged-particles in the earth's magnetotail has been under extensive investigation for the past two decades in order to fully understand the properties of the magnetotail plasmas from both the experimental and theoretical points of view. A number of specialized aspects of the single-particle motion have been analyzed using various methods including approximate analytical methods (Speiser, 1965; Alexeev and Kropotkin, 1970; Sonnerup, 1971; Stern and Palmadesso, 1975; Cowley, 1978; Pellat and Schmidt, 1979) and numerical methods (Speiser, 1967; Cowley, 1971; Eastwood, 1972; Swift, 1977; Wagner et al., 1979; Gray and Lee, 1982; Speiser and Lyons, 1984).

The magnetotail field may be modeled, in its simplest form, by a neutral sheet magnetic profile $B_o(z)\hat{\underline{x}}$ with $B_o(z = 0) = 0$ and a superimposed component $B_n\hat{\underline{z}}$ normal to the plane of the neutral sheet (the so-called "quasi-neutral sheet" geometry). We use the standard magnetotail coordinate system with x in the earthward direction. In the above works, the fundamental question regarding the integrability of the particle motion was not addressed. However, it has recently been shown (Chen and Palmadesso, 1984a) that the magnetotail-like system is intrinsically nonintegrable due to the presence of the normal component B_n. Subsequently, Chen and Palmadesso (1985), henceforth referred to as C-P, have identified various phase space structures in detail. More specifically, using the Poincaré surface of section method, they have shown that particle motions in the magnetotail-like configuration can be classified into three distinct types of orbits occupying disjoint regions of the phase space; the bounded integrable orbits, unbounded stochastic orbits and unbounded transient orbits.

In studying individual orbits, a number of researchers have noted "randomness" in certain orbits. For example, Swift (1977) noted that some orbits appear to randomize in the equatorial plane after a few crossings. Wagner et al., (1979) found that certain orbits exhibit sensitive

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dependence on initial conditions. Gray and Lee (1982) noted that the magnetic moments for some particles exhibit randomness across the equatorial plane. We now understand that this randomness is a manifestation of the intrinsic nonintegrable nature of the system. The work of Chen and Palmadesso (1984a; 1985) provides a systematic and unified understanding of the nature and properties of the particle motion throughout the entire phase space. More recently, Martin (1985) showed that particle motion near an X-point is chaotic. Also, it has long been known (Dragt and Finn, 1976) that motion in the dipole field is not integrable.

The integrability of the particle motion is not purely of mathematical interest. The existence of distinct classes of orbits has profound effects on the plasma particle distribution and on the response of the magnetotail plasmas to external influences. In this regard, the concept of "differential memory" was introduced by C-P to describe the property that particles in the disjoint regions of the phase space respond to external influences on different time scales. This implies that a plasma distribution function has a natural tendency to develop non-Maxwellian features in response to changes in physical conditions. In turn, the distribution and character of the charged-particle orbits can have a significant impact on other dynamical properties of the magnetotail. One important example is the collisionless tearing-mode instability (Furth, 1962; Pfirsch, 1962; Laval et al., 1966) which has long been thought to play an important role in magnetic field reconnection (Coppi et al, 1966; Schindler, 1966). It has recently been shown that the growth rates of the collisionless tearing instability can be enhanced by up to a few orders of magnitude due to the presence of temperature anisotropy or other non-Maxwellian features in the particle distribution (Chen and Palmadesso, 1984b; Chen et al., 1984; Chen and Lee, 1985). Thus, the particle dynamics plays a fundamental role in influencing the magnetotail dynamics. A fundamental question, then, is whether the magnetotail can sustain free-energy carrying non-Maxwellian features that can support large-scale instabilities of potential relevance to magnetotail dynamics, such as the non-Maxwellian collisionless tearing mode.

In this paper, we will examine the way the magnetic field topology of a quasi-neutral sheet can generate non-Maxwellian features in plasma distribution functions and the attendant free-energy that can drive certain

plasma instabilities. In Sec. II, we review the particle dynamics using the Poincaré surface of section method. In Secs. III and IV, physical implications for magnetospheric dynamics will be discussed. We will primarily refer to the ion motion for illustration. The modifications needed for electrons are trivial.

II. Single-Particle Dynamics

In order to model the motion of charged particles in the magnetotail, we consider a magnetic field given by $\underline{B} = B_o f(z)\hat{\underline{x}} + B_n\hat{\underline{z}}$, where B_o is the asymptotic field in the x direction, B_n is the uniform normal field and $B_o f(z)$ is a neutral sheet profile such that $f(-z) = -f(z)$. We will primarily use the Harris configuration

$$f(z) = \tanh(z/\delta), \tag{1}$$

where δ is the characteristic scale length of the magnetic field. The treatment is also applicable to other quasi-neutral sheet configurations and another example, $f(z) = z/\delta$, is discussed in C-P.

We choose the gauge such that the vector potential is $A_y(x,z) = -B_o F(z) + B_n x$, where $dF(z)/dz = f(z)$. The single-particle motion is described by the equation of motion

$$m\frac{d\underline{v}}{dt} = \frac{q}{c}\underline{v} \times \underline{B} \tag{2}$$

This vector equation possesses three exact constants of motion; the Hamiltonian $H = mv^2/2$ where $v^2 = v_x{}^2 + v_y{}^2 + v_z{}^2$, the canonical momentum $P_y = mv_y + (q/c)A_y(x,z)$ and a constant $C_x = m(v_x - \Omega_n y)$ associated with the x-motion. Here, we use $\Omega_n = qB_n/mc$ and $\Omega_o = qB_o/mc$ for each species. If we take the Poisson brackets of the constants of motion, we find $[H, P_y] = 0$ and $[H, C_x] = 0$. However, for C_x and P_y, we find

$$[C_x, P_y] = - m\Omega_n \tag{3}$$

so that C_x and P_y are not in involution. This indicates that the particle orbits may be stochastic. As a general remark, a Hamiltonian system with N degrees of freedom is integrable if and only if there exist N constants of motion that are in involution. Physically, the existence of such global

constants of motion means that one can find a canonical transformation to a frame in which the N coordinates are cyclic. The present system is an interesting example of a nonintegrable system which possesses three exact constants. For more detailed properties of equation (2), see C-P.

In describing the orbits, the following normalization will be used: $b_n = B_n/B_o$, $X = (x - P_y/m\Omega_n)/b_n\delta$, $Y = (y + C_x/m\Omega_n)/b_n\delta$, $Z = z/b_n\delta$, $\tau = \Omega_n t$, and $\hat{H} = H/(mb_n^2\Omega_n^2\delta^2)$. A useful technique for displaying the long-time properties of the particle motion is to use the Poincaré surface of section method (see, for example, Lichtenberg and Lieberman, 1983). For our purpose, a surface of section plot at $Z = 0$ for a given value of $\hat{H}$ is constructed by following an orbit by numerical integration and recording the coordinates X and $\dot{X} = dX/d\tau$ at each point where the orbit crosses the equatorial plane. Figure 1(a), reproduced from C-P, shows such a plot for $\hat{H} = 500$ and $b_n = 0.1$, evaluated at $Z = 0$. All kinematically allowable orbits are confined within the circle of radius $(2\hat{H})^{1/2}$.

The orbits in the region marked A are bounded and integrable. There exists an additional invariant in this localized region of the phase space. This means that the motion of an integrable orbit is constrained to a two-dimensional invariant surface whose cross-section through $Z = 0$ is a closed curve. However, this additional invariant is not a global isolating constant and is not expressible in closed form in terms of elementary functions. Note also that integrable orbits are not necessary adiabatic.

The figure also clearly demonstrates that there is a large stochastic region, marked B. The stochastic region is disjoint from the integrable region A; there is no orbit that can connect the two regions. The orbits in region B are stochastic in the sense that orbits are sensitively dependent on initial conditions, with two nearby orbits diverging rapidly with time. Far from the equatorial plane ($z >> \delta$), the motion is regular. We note that Fig. 3 of Wagner et al. (1979) corresponds to a stochastic orbit.

The nonintegrable orbits may be thought of as forming two flux tubes, which are mirror images of each other, originating from and escaping to infinity. The regions C1 through C5 are the $Z = 0$ cross-sections of the flux tubes as they thread the equatorial plane. All orbits at infinity that can reach $Z = 0$ are mapped into region C1. The orbits then successively cross regions C2 through C5. These regions, C1 through C5, have interesting substructures as shown in Fig. 1(b). A subset of the

orbits entering C1, those crossing S1 and T1, enter the stochastic region B after crossing S5 or T5. That is, region B can only be accessed from region C5. The remainder of orbits escape to infinity after crossing C5 or C4 just above T4. The latter orbits are referred to as transient orbits by C-P. These transient orbits appear to be the type studied by Speiser (1967). Only orbits of this type are shown in Figure 1(b). In Fig. 2, we show a surface of section plot for $\hat{H} = 500$ and $b_n = 0.15$. Note that all the basic features are similar to those in Fig. 1(a) ($b_n = 0.1$). However, with a stronger B_n component, a larger fraction of the phase (i.e., the regions corresponding to C1 through C5) is occupied by the flux tube structure and the transient orbits therein. This is because the field lines are "straighter" for larger b_n so that the orbits can translate in the z direction more easily. (In the limit $b_n \to \infty$, all orbits move in the z direction freely.) Similar structures exist in the surfaces of section for different values of $\hat{H}$. In general, as the value of $\hat{H}$ is reduced, large-scale integrable regions become more fragmented and complicated. For $b_n = 0.1$, large integrable regions vanish for $\hat{H} \leq 6.2$ and essentially all orbits are of the stochastic or the transient types. For the parabolic case, similar features exist. A more detailed description of the phase space structures for both cases can be found in C-P.

III. Current Distributions

The surface of section plots provide another important piece of information. It is clear that the integrable (regular) orbits in region A carry no net current. On the other hand, the nonintegrable orbits, i.e., the transient and stochastic orbits, do carry currents because they come from and escape to infinity. For example, for $\hat{H} = 500$ in the Harris case, all nonintegrable orbits (ions) enter the equatorial plane through region C1 (Fig. 1(a)) and escape from C5 and the nearby hashed regions. This means that regardless of the actual path for different orbits, there is a net drift in the +Y direction in the vicinity of the equatorial plane ($|Z| \leq \delta$). The electron motion is, of course, the opposite so that the electron and ion currents are additive. For the parabolic case in which essentially all nonintegrable orbits are eventually reflected back, the orbits still set up a net drift near the equatorial plane. In this case, the return drift is established away from the $Z = 0$ plane with zero total drift integrated over all Z. Thus, for both the Harris and parabolic case, the non-integrable orbits form a net current in the dawn-to-dusk

direction near the equatorial plane. Note that Stern and Palmadesso (1975) showed that particles which are trapped between mirror points exhibit no net drift across the equatorial plane. Subsequently, Cowley (1978) and Pellat and Schmidt (1979) generalized the proof. From our analysis, we conclude that the proof is exactly applicable to the integrable orbits. The transient and stochastic orbits in the Harris case do carry a net drift because the motion is unbounded in the z direction. In the parabolic field, the proof is still applicable but there is a net drift near the $Z = 0$ plane as described above.

IV. Generation of Free-Energy

A prominent feature of the particle motion in the quasi-neutral sheet geometry is the existence of disjoint regions in the phase space near $Z = 0$, each region consisting of orbits of distinct nature; (1) bounded integrable orbits, (2) unbounded stochastic orbits and (3) unbounded transient orbits. In effect, there exist boundary surfaces between the disjoint regions which do not allow randomization of information or particle energy because of the phenomenon of differential memory (Chen and Palmadesso, 1985). If noise fields are introduced to the magnetic field, the boundary surfaces will break up and allow "diffusion" of orbits. The time scale for this process depends on the strength of the noise field and is slow for low levels of noise. The basic disjointness of the regions and differential memory are expected to be insensitive to low-level noises.

The concept of differential memory has a number of physically interesting implications. Suppose the system contains a population of charged particles in thermal equilibrium. There must also be a steady supply of particles in the distant regions to maintain equilibrium. If the parameters of the distant plasma distribution are changed, then the fact the orbits are divided into distinct types belonging to disjoint regions of the phase space results in a highly non-Maxwellian distribution because of differential memory. The resulting non-Maxwellian distributions may be written in the form $F(H,P_y)$ for the nonintegrable regions and $F(H,P_y,g)$ for the integrable regions where g is the additional invariant of motion. The existence of boundary surfaces can thus establish free-energy which can be tapped by suitable plasma instabilities.

In the context of the earth's magnetotail, it is believed that the collisionless tearing instability may play an important role in

reconnection processes. It has recently been shown (Chen and Palmadesso, 1984b; Chen and Lee, 1985) that the growth rate of the collisionless tearing mode in a non-Maxwellian neutral sheet can be substantially greater, by a few orders of magnitude, than in the Maxwellian case. The general form of the distribution functions which can provide the necessary free-energy is $F(H,P_y,C)$ where C is an independent constant of motion. Thus, if the magnetotail which is initially in thermal equilibrium is subjected to changes in external conditions, e.g., the solar wind, pressure distribution, magnetopause, etc., then the magnetic topology can develop appropriate plasma distributions to drive the non-Maxwellian tearing mode. An important point is that the free-energy for this instability is the particle energy and that the conventional tearing mode also occurs but on faster time scales. That is, changes in external conditions can generate by the process of differential memory an additional source of free-energy in the form of non-Maxwellian distributions.

We add, however, that the above discussion does not take into account the influences of the normal magnetic field on the instability itself under certain circumstances (Galeev and Zelenyi, 1976; Lembege and Pellat, 1982; Coroniti, 1980). This point warrants further consideration and will be addressed in a separate paper. Finally, the above results and discussions are based on single-particle motion. In order to determine more quantitatively the nature and associated time scales of these processes, it is necessary to follow an ensemble of particles. This work is currently underway.

Acknowledgments

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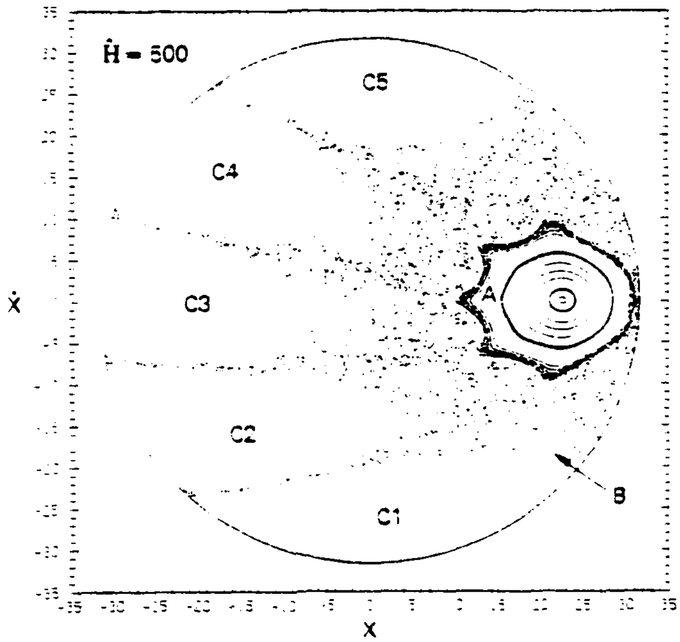

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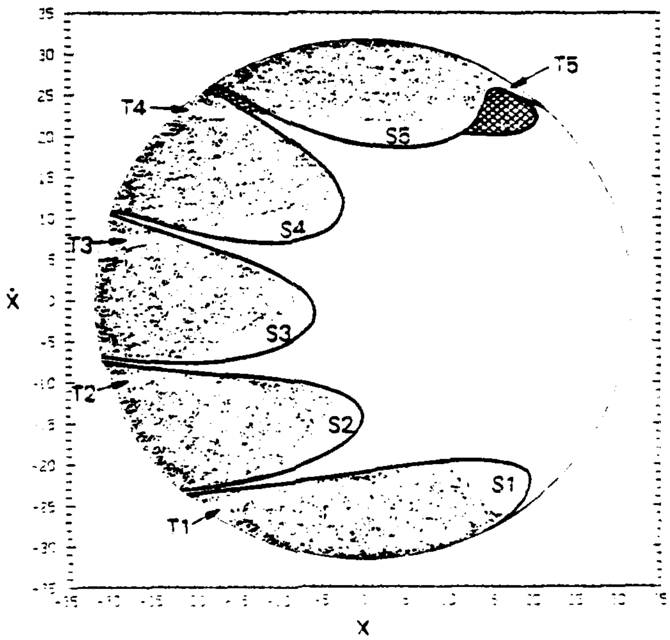

(b)

Fig. 1. (From Chen and Palmadesso, 1985) Surface of section plots for the Harris-type field with $b_n = 0.1$ and $\hat{H} = 500$. (a)Representative integrable orbits in region A and stochastic orbits in region B. 60000 points. (b)Transient orbits, showing the substructures in regions C1 through C5. 42000 points.

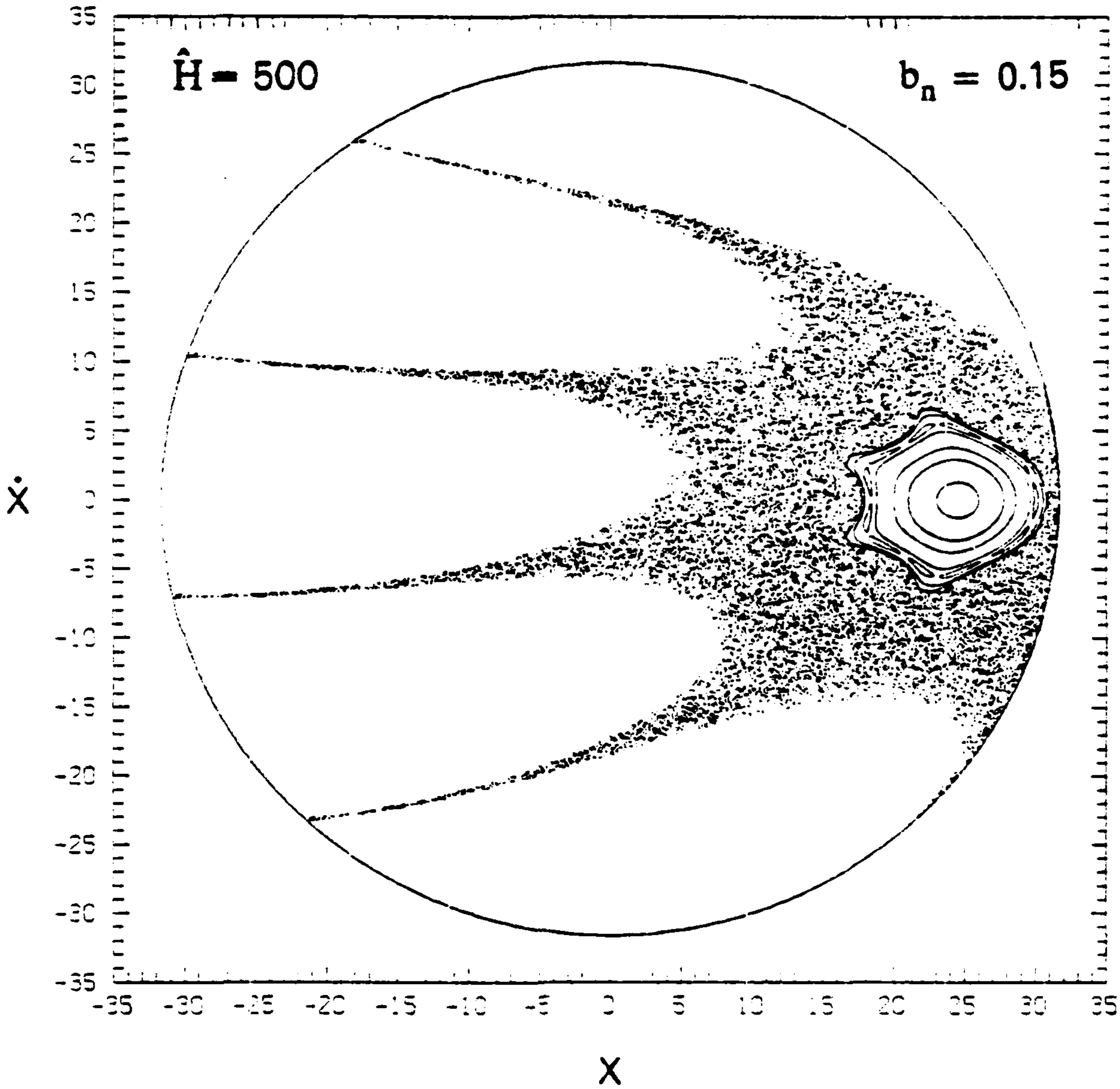

Fig. 2. Surface of section plot. The Harris-type field with $b_n = 0.15$ and $\hat{H} = 500$. The structures are similar to those in Fig. 1.

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